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## DIBARYONS IN NEUTRON STARS

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### ABSTRACT

We investigate the effects of H-dibaryons on the structure of neutron stars. We find that H-particles could be present in neutron stars for a wide range of dibaryon masses. The appearance of dibaryons softens the equation of state, lowers the maximum neutron star mass, and affects the transport properties of dense matter. We constrain the parameter space for dibaryons by requiring that a  $1.44 M_{\odot}$  neutron star be gravitationally stable.

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The existence of long-lived (or even stable) six-quark states with baryon number 2 was first proposed by Jaffe<sup>1</sup> on the basis of bag model considerations; experimental searches for them are being actively pursued,<sup>2</sup> with one recent report of a positive detection.<sup>3</sup> Of the possible dibaryon states, the H-particle, with the quark content of two  $\Lambda$ 's ( $uuddss$ ) in a flavor-singlet spin-zero state ( $J^P = 0^+$ ), is the best candidate for metastability. Bag, Skyrme, quark cluster, and lattice QCD model calculations predict a mass for the H-dibaryon in the range  $m_H \simeq 2.1 - 2.24$  GeV, with most models indicating a value 20 – 80 MeV below that of two  $\Lambda$ 's ( $2m_\Lambda = 2.231$  GeV).<sup>4</sup> In this case, the H-particle would be stable with respect to the strong force, decaying only via weak interactions.

Dibaryons will be present in neutron stars whenever the chemical potential of neutron star matter is greater than the effective mass of the dibaryon. Since they are bosons, like pions and kaons, dibaryons can, in principle, condense in the cores of neutron stars, but with one important difference: dibaryons carry baryon number. If dibaryon interactions are neglected, a Bose condensate of dibaryons would cause the collapse of neutron stars with central densities higher than the minimum density for the appearance of dibaryons.<sup>5</sup> At some level, however, such interactions must be present, and such an instability will not be realized. At very short distances, the interactions among the underlying quark components of the dibaryons will give rise to an effective repulsive interaction between dibaryons. Even a very small self-coupling can stabilize the system against collapse. If the interactions are relatively weak, however, the equation of state of dense matter will be softened by the presence of dibaryons.

At a critical baryon number density,  $n_B \geq n_B^*$ , dibaryons begin to appear inside neutron stars. We study this transition by considering a mixture of neutrons and dibaryons in chemical equilibrium. (For self-consistency, we should also include protons,  $\Lambda$ s, and other light hyperons, but this does not qualitatively change the results for the parameter range of interest.<sup>6</sup>) In equilibrium, the chemical potentials of dibaryons ( $\mu_H$ ) and neutrons ( $\mu_n$ )

are related by

$$\mu_H = 2\mu_n. \quad (1)$$

Of the many candidate equations of state for neutron matter, we choose the medium-stiffness Bethe-Johnson<sup>7</sup> (BJ1) model as a representative and analytically tractable example. (Clearly, it would be of interest to make a systematic study with different nuclear equations of state; here, we have chosen a familiar model to qualitatively illustrate the effects of dibaryons.) In its polytropic incarnation, we can write the energy per neutron as

$$\epsilon_n = \frac{\rho_n}{n_n} = m_n + 236n_n^a \text{ MeV}, \quad (2)$$

where the neutron number density  $n_n$  is given in  $\text{fm}^{-3}$ ,  $m_n$  is the neutron mass in MeV, and  $a = 1.54$ . The neutron pressure and chemical potential can be written as:

$$p_n = n_n^2 \frac{d\epsilon_n}{dn_n} = 363.44n_n^{a+1} \text{ MeV fm}^{-3}, \quad (3)$$

$$\mu_n = m_n + 599.44n_n^a \text{ MeV}. \quad (4)$$

The adiabatic index is then  $\Gamma_n = a + 1 = 2.54$ .

We shall assume that the H-dibaryon is spatially compact, with a size comparable to that of nucleons. This assumption is borne out by bag model estimates<sup>4</sup> and corroborated by a recent study<sup>8</sup> of the quark structure of the H-particle. Then, at the nucleon level, in nuclear mean field theory, the H-particle can be modelled as a complex scalar field,  $\phi$ , with effective Lagrangian<sup>9</sup>:

$$L_\phi = -\frac{1}{2}g^{\mu\nu}\partial_\mu\phi\partial_\nu\phi^\dagger - \frac{m_H^2}{2}\phi^\dagger\phi - \frac{\lambda}{4}(\phi^\dagger\phi)^2 - V[(\phi^\dagger\phi)^n] - L_{int}, \quad (5)$$

where  $n \geq 3$  describes higher order multi-body interactions, and  $L_{int}$  includes the coupling to other hadrons. In mean field theory, the primary effect of the  $H$  interaction with nucleons is to induce a shift  $m_H \rightarrow m_H^*$  in the effective H-mass, due to the mean field of the neutrons. Although density dependent, to first approximation we can absorb this effect in a mean renormalized mass for H; for simplicity, we will denote this effective

mass by  $m_H$  hereafter. As a first approach to modelling the dibaryons that is simple but qualitatively reliable, we will neglect the higher order self-interactions between dibaryons. (These terms may be important at high densities, but in that regime the description of the EOS in terms of nucleons is suspect at any rate.)

We can estimate the dibaryon self-coupling  $\lambda$  by using a variant of the P-matrix formalism of Jaffe and Low,<sup>10</sup> applied to diquarks by Donoghue and Sateesh.<sup>11</sup> If we consider two dibaryons brought together, the energy of the resulting configuration can be written as  $E_2 = 2m_H + \Delta E(r, \lambda)$ , where  $r$  is their separation. As  $r \rightarrow R_4$ ,  $E_2 \rightarrow M_4$ , where  $M_4$  and  $R_4$  are the mass and radius of the 12-quark baryon number-4 state. The mass  $M_4$  can be estimated from the bag model,<sup>4</sup>  $M_4 \simeq 4.7$  GeV. Calculating the interaction energy for the two-dibaryon state confined inside a cavity of radius  $R$ , we find that<sup>11</sup>  $E_{int}(R, \lambda) = 0.34\lambda/4E^2 R^3$ , where  $E^2 = p^2 + m_H^2$  and the momentum  $p = \pi/R$  from the vanishing of  $\phi$  at the cavity boundary. Taking the cavity to be the size of the baryon number-4 state,  $R = R_4$ , and equating  $\Delta E(R_4, \lambda) = E_{int}(R_4, \lambda)$ , we obtain

$$\lambda = 11.8R_4^3(M_4 - 2m_H) \left( \frac{\pi^2}{R_4^2} + m_H^2 \right). \quad (6)$$

For the range  $R_4 = 1.1 - 1.4$  fm and  $m_H = 2 - 2.2$  GeV, we estimate  $\lambda \simeq (3.4 - 13) \times 10^3$ . Given the crudeness of the approximations involved, this estimate should be considered a guide to the expected order of magnitude for the coupling.

Since the self-coupling  $\lambda$  lies comfortably above a critical value  $\lambda^* = 4\pi m_H^2/m_{Pl}^2 \simeq 4 \times 10^{-37}$ , we can use the equation of state for self-interacting bosons derived by Colpi, Shapiro, and Wasserman<sup>12</sup> to excellent approximation (corrections are of order  $\mathcal{O}(\lambda^*/\lambda) \sim 10^{-40}$ ). For a spherically symmetric scalar field with time-independent energy-momentum tensor, the lowest energy solution is of the form

$$\phi(r, t) = \Phi(r)e^{-i\omega t}, \quad (7)$$

where  $\Phi(r)$  is a real function of the radial coordinate. In the limit  $\lambda \gg \lambda^*$ , the solution

to the scalar field equation of motion to lowest order in  $\lambda^*/\lambda$  is just

$$\Phi = \frac{m}{\lambda^{1/2}} \left( \frac{\omega^2}{m^2} - 1 \right)^{1/2} \quad \text{for } r < R, \quad (8)$$

where  $R$  denotes the radius of the configuration.<sup>13</sup> The dibaryon equation of state is obtained by substituting this solution into the scalar field energy-momentum tensor,

$$T_\nu^\mu = \frac{1}{2} g^{\mu\sigma} \left( \partial_\sigma \phi^\dagger \partial_\nu \phi + \partial_\sigma \phi \partial_\nu \phi^\dagger \right) - \frac{1}{2} \left[ g^{\lambda\sigma} \partial_\lambda \phi^\dagger \partial_\sigma \phi + m_H^2 \phi^\dagger \phi + \frac{1}{2} \lambda (\phi^\dagger \phi)^2 \right] \quad (9)$$

and the H-particle current

$$J^\mu = \frac{i}{2} g^{\mu\nu} \left( \phi^\dagger \partial_\nu \phi - \phi \partial_\nu \phi^\dagger \right). \quad (10)$$

Defining the dimensionless parameter  $z = \omega^2/m^2$ , the dibaryon energy density and pressure are then given by:

$$\rho_H = \rho_0(z-1)(3z+1), \quad (11)$$

$$p_H = \rho_0(z-1)^2, \quad (12)$$

where

$$\rho_0 = m_H^4/4\lambda = 523 \frac{\text{MeV}}{\text{fm}^3} \left( \frac{m_H}{2000 \text{ MeV}} \right)^4 \left( \frac{10^3}{\lambda} \right). \quad (13)$$

The parameter  $z$  is related to the baryon number density of dibaryons by:

$$n_B = 2n_H = \frac{8\rho_0}{m_H} (z-1)\sqrt{z}. \quad (14)$$

(Thus,  $z \rightarrow 1$  as the dibaryon density goes to zero.) This specifies the dibaryon equation of state in parametrized form. (Alternatively, eliminating  $z$ , the equation of state can be written<sup>12</sup>  $p_H = (4\rho_0/9)[(1 + 3\rho_H/4\rho_0)^{1/2} - 1]^2$ .) We can then calculate the dibaryon chemical potential:

$$\mu_H = \frac{\rho_H + p_H}{n_H} = m_H \sqrt{z}, \quad (15)$$

and the adiabatic index

$$\Gamma_H = \frac{p_H + \rho_H}{p_H} \left( \frac{dp_H}{d\rho_H} \right) = \frac{4z}{3z-1}. \quad (16)$$

At low density ( $n_H \rightarrow 0$ ),  $\Gamma_H \rightarrow 2$ , and at high density ( $\rho \gg \rho_0$ )  $\Gamma_H \rightarrow 4/3$ ; the latter reflects the fact that the dibaryon EOS reduces to that of a free relativistic gas in the high density limit. Since  $\Gamma_H < \Gamma_n$ , dibaryons generally soften the equation of state.

Imposing chemical equilibrium between neutrons and dibaryons, we find that the threshold baryon number density for the appearance of dibaryons is:

$$n_B^* = \left( \frac{m_H - 2m_n}{1198.88 \text{ MeV}} \right)^{1/a} \text{ fm}^{-3} = 0.45 \text{ fm}^{-3} \left( \frac{m_H - 2m_n}{2m_\Lambda - 2m_n} \right)^{1/a}. \quad (17)$$

This threshold is independent of the self-coupling  $\lambda$ . For  $m_H \simeq 2m_\Lambda$ , the threshold baryon density for the appearance of H-particles is about three times nuclear matter density.

For densities above  $n_B^*$ , the equation of state will be that of a neutron-dibaryon mixture:

$$\begin{aligned} n_B &= n_n + 2n_H, \\ \rho &= \rho_n + \rho_H, \\ p &= p_n + p_H = p(n_B), \\ \mu_B &= (\rho + p)/n_B = \mu_B(p). \end{aligned} \quad (18)$$

Imposing chemical equilibrium, the system can be described with one parameter, for example the neutron density  $n_n$ .

We have constructed models of neutron stars by integrating the Oppenheimer-Volkoff equation of hydrostatic equilibrium using the above equation of state. The effect of the appearance of dibaryons can be seen in Fig. 1, which shows the neutron star mass  $M$  vs. radius  $R$  for models with (solid curve) and without dibaryons (dashed curve). For this choice of parameters,  $m_H = 2200 \text{ MeV}$  and  $\lambda = 8000$ , the maximum neutron star mass is reduced when dibaryons are included, a consequence of the softening of the equation of state.

In Fig. 2, we plot the neutron star mass threshold for dibaryons to be present: a neutron star with BJ1 equation of state and a mass above the threshold will have an

admixture of dibaryons in its core. Note that dibaryons are present in stars of mass  $1.44M_{\odot}$  for a wide range of H-particle masses. In particular, even if the H is unstable ( $m_H > 2m_A$ ), it seems to be welcome.

The dependence of the neutron star maximum mass on the dibaryon mass and coupling constant is shown in Fig. 3. The constraint that the maximum neutron star mass exceed  $1.44M_{\odot}$ , the mass of the binary pulsar<sup>14</sup> PSR1913+16, implies a lower bound on the strength of dibaryon interactions ( $\lambda$ ) for a given dibaryon mass. In Fig. 4, we show this constraint on the dibaryon parameter space for the neutron-dibaryon model. Interestingly, this lower bound on  $\lambda$  as a function of  $m_H$  is close to the bag model estimate of  $\lambda$ , Eqn.(6). We also note that if  $m_H < 2m_n = 1879$  MeV, then "dibaryon matter" would be absolutely stable at zero pressure; in this case, dibaryons would form 'boson stars'<sup>12</sup> (perhaps with a small crust of nuclear matter) with maximum mass  $M_{max} = 0.22(\lambda/\lambda^*)^{1/2}m_{\pi}^2/m_H = 1.44M_{\odot}(\lambda/2.8 \times 10^3)^{1/2}(2m_n/m_H)^2$ .

Aside from their effects on neutron star structure, dibaryons may also have important consequences for the transport and cooling properties of neutron stars. Since they are a non-ideal Bose gas, the low-lying excitations will have a phonon spectrum, and the dibaryons will presumably act as a superfluid. In fact, one can demonstrate this directly by considering the spectrum of small fluctuations<sup>15</sup> about the solution (7). Also, since they are not Pauli-blocked, neutrino emission from H-H and H-n scattering may be considerably enhanced over the modified URCA process, thereby accelerating the early stages of neutron star cooling. We conclude that dibaryons could play an important role in neutron star physics.

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## FIGURE CAPTIONS

**Figure 1.** Gravitational mass  $M$  versus stellar radius  $R$  for BJ1 equation of state (solid curve) and for BJ1 with dibaryons of mass  $m_H = 2200 \text{ MeV}$  and self-coupling  $\lambda = 8000$  (dashed curve).

**Figure 2.** Threshold neutron star mass above which dibaryons appear, as a function of dibaryon mass.

**Figure 3.** Maximum neutron star mass as a function of  $\lambda$  for  $m_H = 2000, 2100, 2200, 2300 \text{ MeV}$ . The horizontal line at  $1.44M_\odot$  denotes the mass of the binary pulsar.

**Figure 4.** Minimum value of  $\lambda$  required for stability of a  $1.44M_\odot$  neutron star,  $\lambda_{\min}/1000$ , as a function of H-mass  $m_H$ .

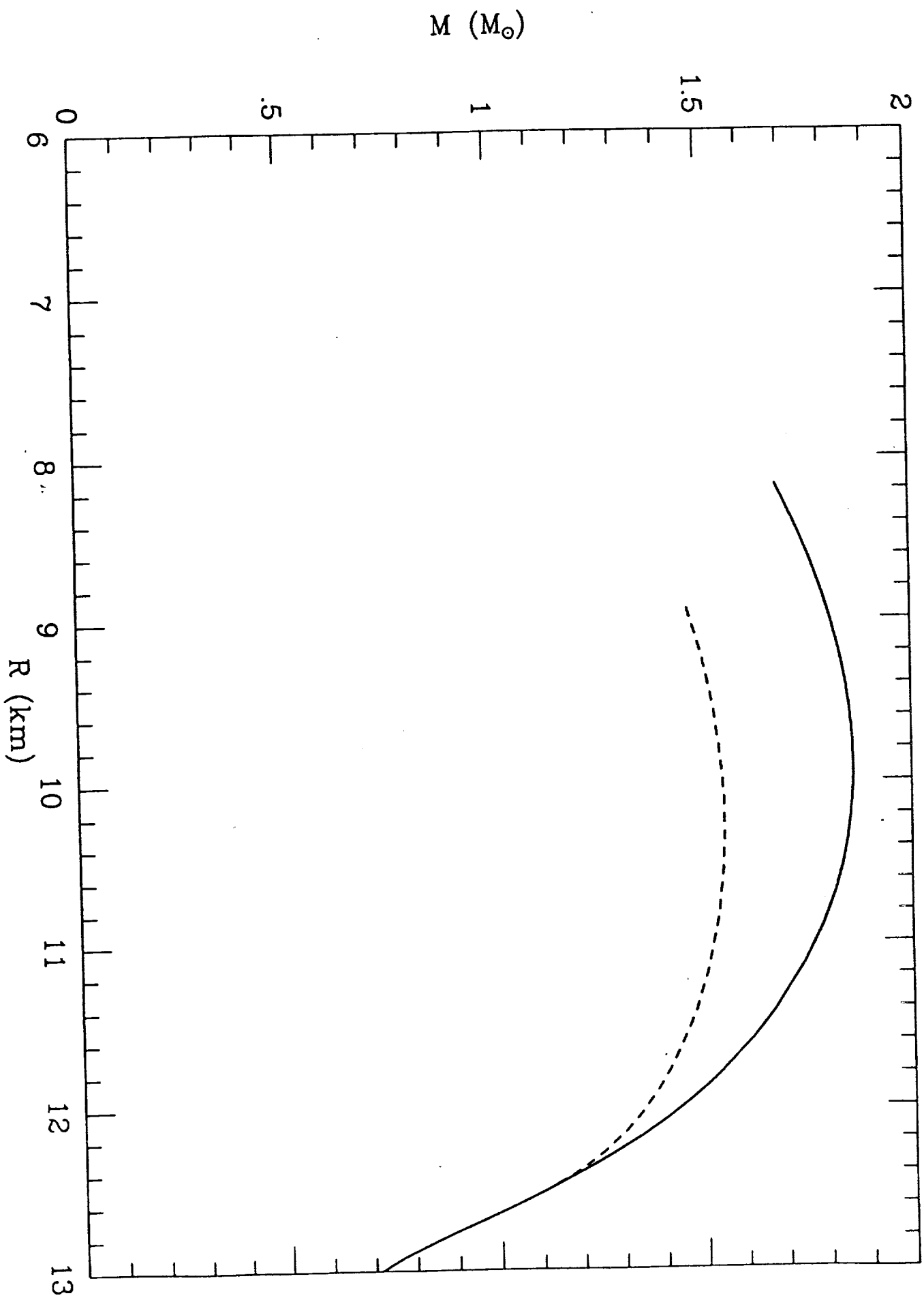


Fig. 1

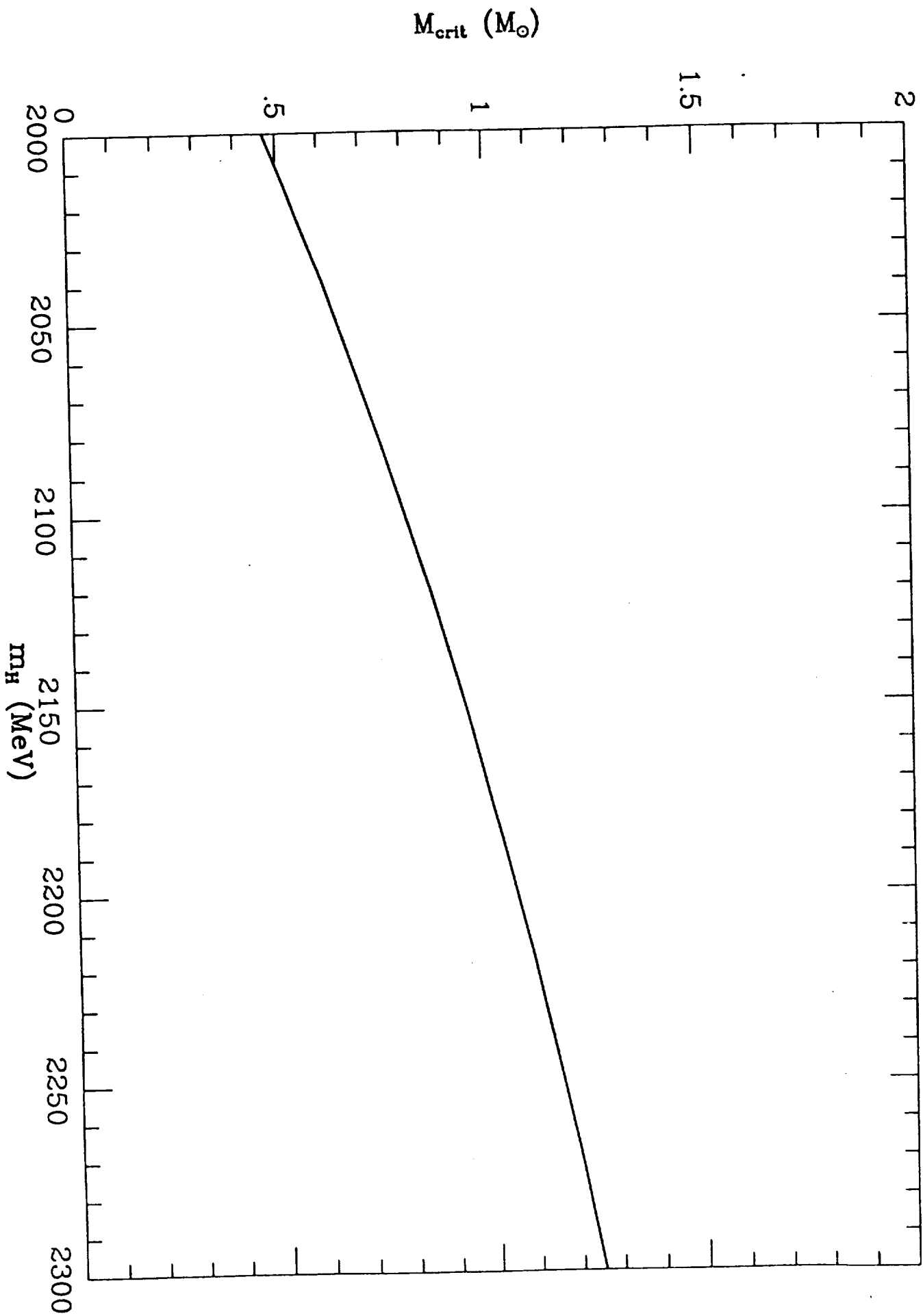


Fig. 2

